An Extra Push from Entrance-Channel Effects

N. Rowley [a] N. Grar [b] K. Hagino [c]

[a] *Institut de Recherches Subatomiques/Université Louis Pasteur (UMR 7500), 23 rue du Loess, F-67037 Strasbourg, France*

[b] *DAC, Department of Physics, Université de Sétif, Algérie*

[c] *Department of Physics, Tohoku University, Sendai 980-8578, Japan*

Abstract

Symmetric heavy-ion collisions are known to display an 'extra-push' effect. That is, the energy at which the s-wave transmission is 0.5 lies significantly higher than the nominal Coulomb barrier. Despite this, however, the capture cross section is still greatly enhanced below the uncoupled barrier. It is shown that this phenomenon can be simply explained in terms of entrance-channel effects which account for long-range Coulomb excitations.

Key words: Coupled channels, Coulomb barrier distribution, Fusion and fusion-fission reactions, Coulomb excitation, Extra push
PACS: 24.10.Eq, 25.70.Hi, 25.70.Jj, 25.70.De

When reactions between intermediate-mass heavy ions lead to non-fissile composite systems, the relationship between the cross sections for capture (passing over or penetrating through the Coulomb barrier), fusion (evolution to a compact equilibrated compound nucleus; CN) and evaporation residues (ER) is straightforward. If fission is unimportant, all of the above cross sections are essentially equal: $\sigma_{\rm cap} = \sigma_{\rm fus} = \sigma_{\rm ER}$. Of course it is well known that couplings to collective states of the target and projectile can lead to a distribution of Coulomb barriers [1] but this does not in any way change the above relationship, any structure in $\sigma_{\rm cap}$ also being present in $\sigma_{\rm ER}$. To study the effects of the entrance channel, one may simply measure the long-lived evaporation residues which recoil in a relatively narrow cone around the beam direction (dispersed by the emission of neutrons, protons and α-particles from the CN). The results for intermediate-mass systems almost invariably show that collective couplings increase the sub-barrier capture cross section (see, for example, Ref. [1]).

* Corresponding author
Email address: neil.rowley@ires.in2p3.fr (N. Rowley).

For heavier systems, other reactions mechanisms intervene and complicate the situation both experimentally and theoretically. In particular, the composite system might not fuse but instead quickly separate into two fragments similar in mass and charge to the target and projectile (quasifission; QF). The CN itself may also fission (fusion-fission; FF) rather than decaying to a long-lived residue through particle evaporation. For very heavy systems the fission modes dominate and a complete understanding of the interplay between the various reaction mechanisms is especially important in heavy-element creation.

To measure $\sigma_{\rm cap}$ directly in the general case, $\sigma_{\rm ER}$, $\sigma_{\rm QF}$ and $\sigma_{\rm FF}$ must all be measured (including the fragment angular distributions) in order to obtain $\sigma_{\rm cap}$. Though if quasifission is not thought to be important, one could still try to obtain the capture cross section by measuring only the evaporation residues, and using an evaporation-model code that accounts for the competition between fusion-fission and fusion-evaporation decay modes to reconstruct the capture cross section required to reproduce $\sigma_{\rm ER}$. This was the aim of a series of experiments performed at GSI using projectiles and targets around mass 100 [2,3,4,5]. The interesting result is that the capture cross sections obtained displayed a so-called extra-push effect. That is, the energy $\bar{B}$ at which the deduced s-wave transmission T_0 was 0.5, could greatly exceed the barrier height predicted by potential models such as that of Bass [6]. This in itself might be explained by an internal barrier which must be crossed after passing the outer Coulomb barrier if fusion is to take place, and this could be thought of as the conditional saddle point in the liquid-drop nuclear potential [7]. However, the data are not entirely consistent with such a description since, despite the shift of the $T_0 = 0.5$ point to higher energies, $\sigma_{\rm cap}$ was still found, as for lighter systems, to be strongly enhanced at energies well below the Bass barrier. This enhancement was quantified by defining a single (adiabatic) barrier $B_{\rm ad}$ which yielded the correct cross section at the very lowest energies, and thus obtaining an overall width of the barrier distribution $D_\infty = \bar{B} - B_{\rm ad}$. For the system ^{100}Mo + ^{100}Mo, for example, it was found that $D_\infty \approx 20$ MeV.

The authors of Ref. [2] tried to fit their data with an entrance-channel model using the simplified coupled-channels code CCFUS [8] with couplings to the known quadrupole- and octupole-phonon states of target and projectile. They found that in general such calculations could account for only about one half of D_∞. The main aim of the present paper is to show that more complete coupled-channels calculations are in fact capable of fitting D_∞ rather well, and also yielding the correct shape of the capture cross section (assumed by Quint et al. to arise from a gaussian barrier distribution; see Fig. 2). An important ingredient missing from the earlier calculations will be shown to be the long-range Coulomb couplings which polarise the target and projectile well before the Coulomb barrier is reached. The role of multi-phonon excitations is also important.

The points in Fig. 1 shows both on a logarithmic scale and a linear scale the deduced experimental s-wave transmission as a function of the incident energy $E_{\rm cm}$ for the system ^{100}Mo + ^{100}Mo. They were derived by assuming a gaussian barrier distribution with a centroid B and standard deviation Δ and varying these parameters until the fusion-evaporation-model code HIVAP [9] reproduced the evaporation-residue cross section. The experimental values of T_0 are then obtained through

$$T_0^{\rm exp} = T_0^{\rm theory} \frac{\sigma_{\rm ER}^{\rm exp}}{\sigma_{\rm ER}^{\rm HIVAP}}. \qquad (1)$$

This is a very good way to represent the data, since the quantity T_0 is directly related to the entrance-channel dynamics. However it should be stressed that the experimental T_0 are not true experimental data. They depend not only on B and Δ but also on the parameters entering into the HIVAP calculation. This leads to certain ambiguities for some system, a point to which we shall return later. For the moment we accept these numbers at face value and attempt to fit them with calculations using the program CCFULL [11], again using known phonon states in ^{100}Mo.

This nucleus has strong quadrupole- and octupole phonon states lying at relatively low excitation energies and we shall use the adopted empirical values of these energies and the corresponding deformation parameters: $E(2^+) = 0.536$ MeV, β_2=0.21; $E(3^-) = 1.908$ MeV, β_3=0.17 [12]. The only other parameters entering our calculations are the no-coupling barrier height $B_{\rm nc}$, which we shall vary to fit the data, and the diffusivity of the nuclear potential for which we take a standard vaue of $a = 0.6$ fm.

The dashed curves in Fig. 1 a,b show the no-coupling result, which is seen to greatly underestimate T_0 at low energies. The other curves show calculations including various phonon couplings $[N_{\rm quad}, N_{\rm oct}]$. The symmetry of the present system allows us to use a simple theoretical trick to reduce the number of channels in a given calculation. For example, the calculation with one quadrupole phonon in both target and projectile, along with the mutual excitation can be exactly treated as a two-channel calculation with renormalised couplings. The details of this method will be presented elsewhere [10]. Thus the calculation labelled $[4, 2]$ means two quadrupole- phonon excitations and one octupole excitation in each nucleus along with all possible mutual excitations. It is clearly seen that as the complexity of the coupling increases, the theoretical results converge to the experimental curve both at high energies (see linear scale) and low energies (logarithmic scale). The final calculation $[4, 2]$, however, still slightly underpredicts T_0 at the very lowest energies, and it might be asked why we do not pursue this with a $[4, 4]$ calculation.

The problem here is that the full coupled-channels calculations become numer-

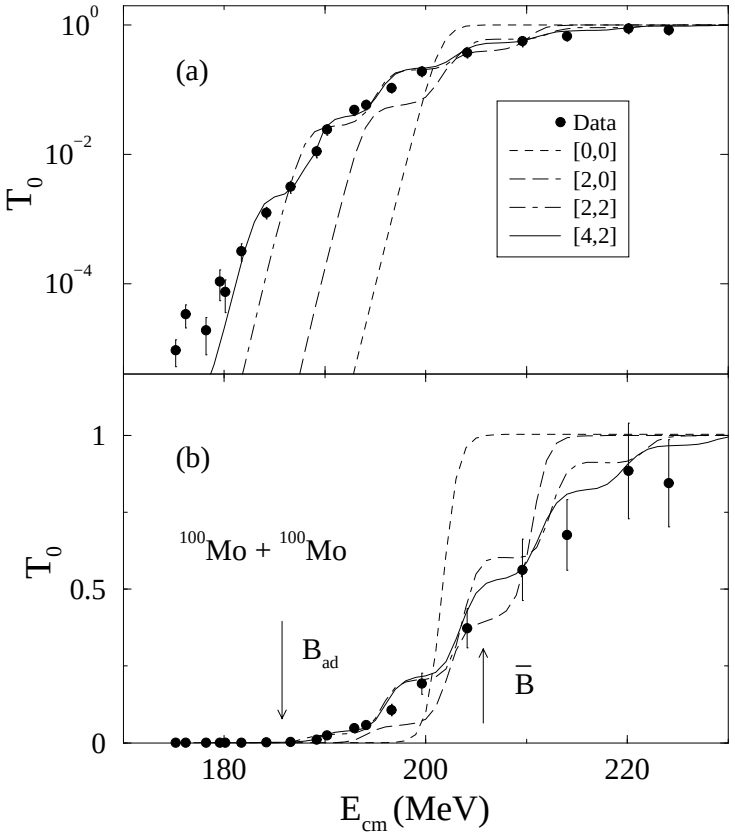

Figure 1. Experimental T_0 compared with various CCFULL calculations with different numbers of phonon excitations. See text for details. Arrows indicate the average barrier $\bar{B}$ and the adiabatic barrier $B_{\rm ad}$, whose difference gives D_∞. Parts (a) and (b) show same curves but on logarithmic and linear scales.

ically unstable at low energies if too many channels are included. The reason is that we are essentially integrating the Schroedinger equation at energies around 30 MeV under the highest effective barrier, and the energies losses due to couplings to the phonon states further reduce the kinetic energy of the relative motion. This problem increases with the number of phonon channels and the program breaks down at the lowest energies. However, the problem may be overcome to some extent by reducing the width of the Coulomb barrier, and this can be achieved by decreasing the diffusivity a. In Fig. 2 we show the results of calculations using $a = 0.2$ fm. We should stress that we do not believe such a low value of the diffusivity but only use it as a means of seeing the effect of the higher phonon couplings in the $[4, 4]$ calculation. However, the use of $a = 0.2$ changes rather little the barrier positions. Its main effect is to decrease the rate at which the cross section falls off below the Coulomb barrier. But since the cross section at low energies is dominated by the lowest barriers, this effect is only significant below the very lowest (adiabatic) barrier.

We show again in Fig. 2a the calculations with the same coupling schemes as in Fig. 1, and note that the inclusion of the double-octupole phonon shifts the low-energy cross section down by about a further 2 MeV. We would, of course, obtain a similar shift with the more physical value of $a = 0.6$ fm in Fig. 1a if

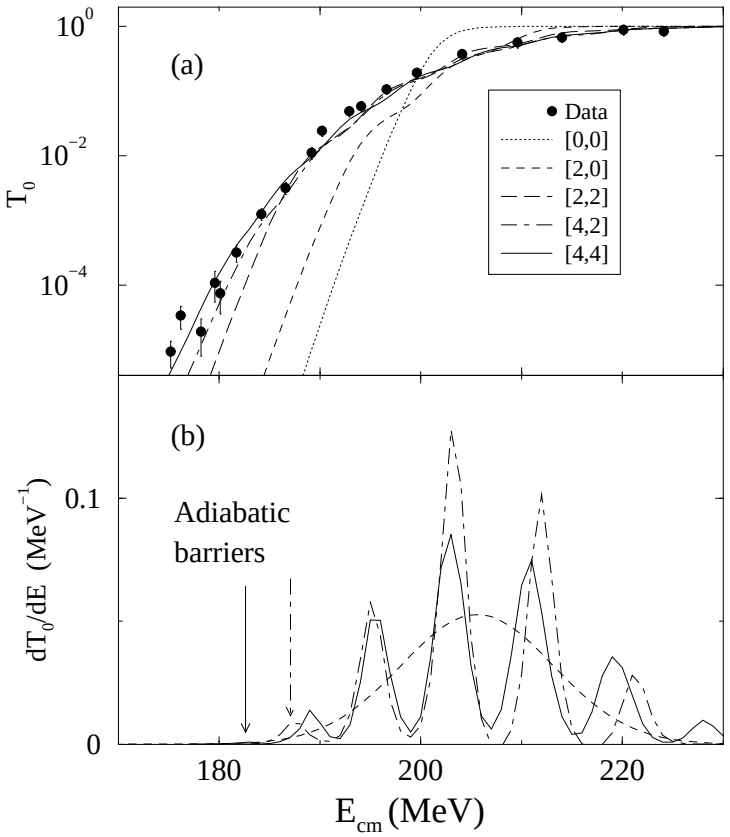

Figure 2. Using $a = 0.2$ fm permits the $[4,4]$ calculation with both double quadrupole- and octupole-phonon excitations. Note that the calculations have virtually converged, with a new lowest barrier emerging but with very small weight. The dashed curve in (b) is the gaussian barrier distribution of Ref. [2].

it were possible to do this calculation. We do not insist too much on this fine detail of the problem since, as already noted, there are ambiguities stemming from the HIVAP calculation. We have also ignored other possible coupling effects such as neutron-transfer channels, though these will always have unfavourable Q values for symmetric systems. Fig. 2b shows the derivative of T_0 with respect to the incident energy for the $[4,2]$ and $[4,4]$ calculations. It is well known that this gives the distribution of barriers $D(E)$ [13], and it can be seen that there is little difference between the two distributions except for the presence of a lower adiabatic barrier with very small weight (barely visible on this scale) in the latter case. We can, therefore, conclude that the calculations have essentially converged. This is reassuring since the need to introduce higher phonon states might be somewhat dubious. We note that the adiabatic barrier of our calculations is not the same as that of Quint et al. which has a weight of 1 and is supposed simply to reproduce T_0 at low E.

The calculations that we have presented show the importance of higher phonon

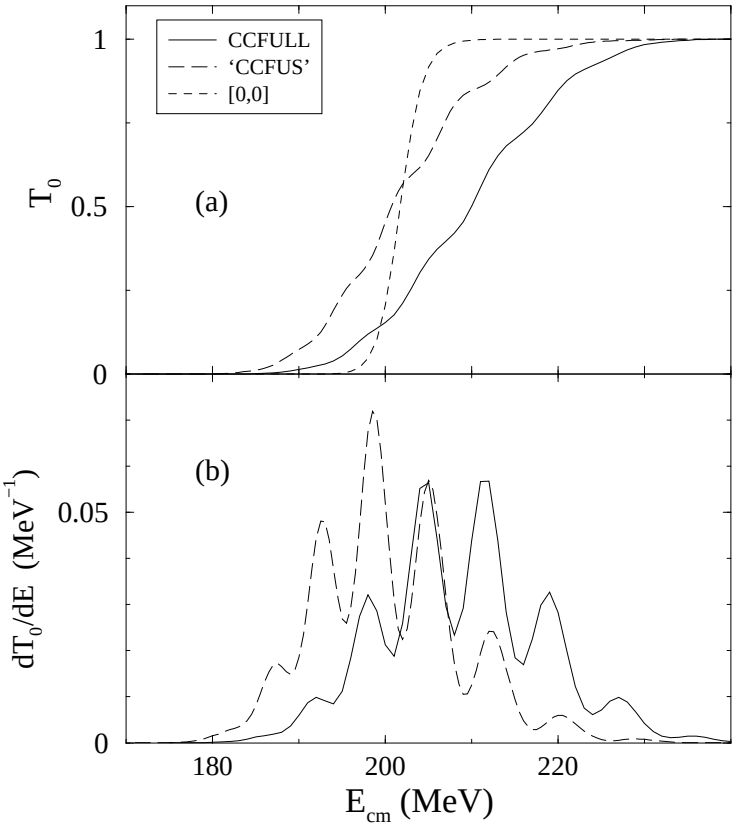

Figure 3. The $[4,4]$ CCFULL calculation compared with a $[4,4]$ calculation in the spirit of CCFUS. See text for details. Note that the latter calculation does not produce a shift of the $T_0 = 0.5$ point, whereas the CCFULL calculation gives a shift of about 10 MeV due to the higher weights of the high-E barriers.

couplings not included in the CCFUS calculations of Ref. [2]. There is, however, another very important difference which introduces new physics into the barrier distribution, and which we shall now elaborate.

In CCFUS, everything is essentially determined in the barrier region, and the barrier heights and weights obtained through the diagonalisation of the coupling matrix (including excitation energies) at the barrier radius. This is probably a reasonable approximation for the short-ranged nuclear field but will fail for heavy systems where the Coulomb field plays an important role at large distances. In order to simulate a CCFUS-type model but still include all of the nuclear $[4,4]$ couplings, we performed a calculation in which the Coulomb deformation parameters were set to zero. However, this will also change the barrier heights, since the deformed Coulomb field is not negligible at the barrier. In order to correct for this, we renormalised the nuclear deformation parameters (this is possible since the same geometrical factors appear in both couplings). The results for the relevant barrier distributions are shown in Fig. 3. One sees that the barriers occur at almost exactly the same positions in the two calculations but that in the complete calculation the weights are greatly shifted towards the high-energy barriers, due to the Coulomb couplings at large distances. In effect, the Coulomb field favours the linear superposition

of states which lowers its own energy. Since it has the opposite sign from the nuclear field, this configuration is precisely that which minimises the nuclear forces, that is, the one corresponding to the *highest* barrier. In other words, the nuclei are polarised in the entrance channel to disfavour the lower barriers. The effect leads to an overall shift of the barrier centroid of around 10 MeV, even though the individual barrier positions remain unchanged. (The $T_0 = 0.5$ point of the CCFUS-type calculation is essentially unshifted.) Since D_∞ in the present case is about 20 MeV, this gives the factor of around 2 which was missing from D_∞ in the calculations of Quint et al.

We believe that similar considerations apply to the work of Berdichevsky et al. [15] who used a single-particle model to approximately derive the barrier splittings but without doing a full calculation of the scattering. (They rather compared their spread of barriers with the Δ of Ref. [2].) Such a model may give a reasonable spread of barriers but it is important to have the relevant correlations which render the nuclear states collective in order to get the correct reaction dynamics and the correct shape of T_0.

We have obtained an excellent fit to the proposed shape of the capture cross section with physically reasonable parameters. However, we should now return to the question of what is the appropriate uncoupled barrier height. Do our calculations retrieve the Bass barrier? The answer to this question is no. Our uncoupled barrier is 201.7 MeV and the Bass barrier 195.2 MeV. That is we still need an uncoupled barrier 6.5 MeV higher than $B_{\rm Bass}$ (previously 12.2 MeV [2]) and we should ask why this is so. There are various possible explanations for this including:

- The Bass potential contains a factor $R_1\,R_2/(R_1 + R_2)$ which accounts for the curvature of the two nuclear surfaces. This factor is largest for symmetric systems and may simply over-estimate the potential for such reactions, giving too low a barrier.
- The Bass potential parameters are fitted to experimental data, which necessarily contain all possible couplings. It is known that high-lying phonon states shift the barrier centroid to lower energies [14]. Thus the uncoupled barrier should probably be taken to be higher than the Bass barrier if one accounts for the couplings explicitly, as we do here.

We should not, however, forget the ambiguities in mapping from $\sigma_{\rm ER}$ to $\sigma_{\rm cap}$. These come both from ambiguities in the statistical-model parameters and from the complete neglect of the QF process, and in this context it is interesting to look at other symmetric systems. Fig. 4 shows our fits to the systems ^{90}Zr + ^{90}Zr [3] and ^{100}Mo + ^{110}Pd [2]. These will be discussed in detail elsewhere [10]. Here we note simply that the barrier shift we require for ^{90}Zr + ^{90}Zr is 4.1 MeV, similar to that for ^{100}Mo + ^{100}Mo, but for ^{100}Mo + ^{110}Pd we require a shift of 15 MeV (previously 29.0 MeV), which does not seem con-

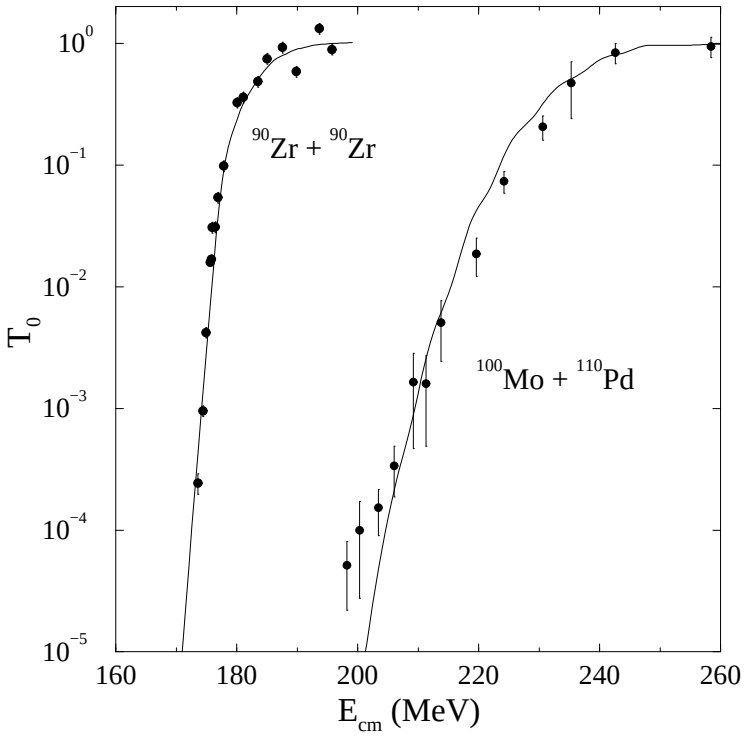

Figure 4. CCFULL fits to two different systems. Large ambiguities exist in the experimental curve for the heavier one. See text.

sistent with the other systems. However, it has been pointed out [2,3] that if one performs the HIVAP calculations with a smaller shell-damping parameter (the energy range over which shell effects are smeared out) different solutions for the gaussian parameters (hence different T_0) are possible. The effects are relatively small for ^{90}Zr + ^{90}Zr and ^{100}Mo + ^{100}Mo, changing Δ rather little but moving $\bar{B}$ down to make our uncoupled barrier rather closer to the Bass value. However, for the system ^{100}Mo + ^{110}Pd (where the ratio $\sigma_{\rm ER}/\sigma_{\rm cap}$ is much smaller and $\sigma_{\rm QF}$ may also be more important) the effect is much larger, giving a shift down of around 8 MeV but still leaving the uncoupled barrier around 7 MeV higher than $B_{\rm Bass}$.

The ambiguities here are sufficiently important to merit further experimental investigation. The most pertinent case is ^{100}Mo + ^{110}Pd, and the ambiguity could be resolved by a direct measurement of $\sigma_{\rm cap}$ for this system, as discussed at the beginning of this Letter. It might, however, be simpler to exploit unitarity and obtain the capture barrier distribution from the large-angle quasielastic flux scattered back from the Coulomb barriers [16,17,18].

This work was supported by the Grant-in-Aid for Scientific Research, Contract No. 16740139, from the Japanese Ministry of Education, Culture, Sports, Science and Technology.

References

[1] M. Dasgupta, D. J. Hinde, N. Rowley and A. M. Stefanini, Annu. Rev. Nucl. Part. Sci. 48 (1998) 401.

[2] A.B. Quint et al., Z. Phys. A 346 (1993) 119.

[3] J.G. Keller et al., Nucl. Phys. A 452 (1986) 173.

[4] K.H. Schmidt and W. Morawek, Rep. Prog. Phys. 54 (1991) 949.

[5] C.-C.Sahm et al., Nucl.Phys. A 441 (1985) 316.

[6] R. Bass, Phys. Rev. Lett. 39 (1977) 265; Nuclear Reactions with Heavy Ions (Springer Verlag, Berlin Heidelberg New York, 1980) 152.

[7] S. Bjornholm, W.J. Swiatecki, Nucl. Phys. A 391 (1982) 471; J.P. Blocki, H. Feldmeier and W.J. Swiatecki, Nucl. Phys. A 459 (1986) 145.

[8] R.A. Broglia, C.H. Dasso, S. Landowne, G. Pollarolo, Phys. Lett. B 133 (1983) 34; C.H. Dasso and S. Landowne, Comp. Phys. Comm. 46 (1987) 187; J. Fernandez-Niello, C.H. Dasso and S. Landowne Comp. Phys. Comm. 54 (1989) 409.

[9] W.N. Reisdorf et al., unpublished.

[10] N. Grar and N. Rowley, in progress.

[11] K. Hagino, N. Rowley and A.T. Kruppa, Comp. Phys. Comm. 123 (1999) 143.

[12] ENSDF, http://www.nndc.bnl.gov/ensdf/

[13] N. Rowley, G.R. Satchler and P.H. Stelson, Phys. Lett. B 254 (1991) 25; A.B. Balantekin and N. Takigawa, Rev. Mod. Phys. 70 (1998) 77.

[14] K. Hagino, N. Takigawa, A.B. Balantekin, Phys. Rev. C 56 (1997) 2104; K. Hagino, N. Takigawa, M. Dasgupta, D.J. Hinde, and J.R. Leigh, Phys. Rev. Lett. 79 (1997) 2014.

[15] D. Berdichevsky, A. Lukasiak, W. Norenberg, P. Rozmej, Nucl.Phys. A 499 (1989) 609.

[16] H. Timmers et al., Nucl. Phys. A 584 (1995) 190.

[17] K. Hagino and N. Rowley, Phys. Rev. C 69 (2004) 054610.

[18] E. Piasecki et al., Phys. Rev. C 65 (2002) 054611; Phys. Lett. B 615 (2005) 55.